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HAGEDORN BEHAVIOR OF LITTLE STRING THEORIES[1]

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We examine the Hagedorn behavior of little string theory using its conjectured duality with near-horizon NS5-branes. In particular, by studying the string-corrected NS5-brane supergravity solution, it is shown that tree-level corrections to the temperature vanish, while the leading one-loop string correction generates the correct temperature dependence of the entropy near the Hagedorn temperature. Finally, the Hagedorn behavior of ODp-brane theories, which are deformed versions of little string theory, is considered via their supergravity duals.

1 Introduction

The near-horizon limit of NS5-branes is conjectured to be dual to Little String Theory (LST) with 16 supercharges [1]. LST is a 5+1 dimensional non-gravitational and non-local theory of strings [2] (see also [3]). As for any string theory, the statistical mechanics description of LST breaks down at a certain temperature, known as the Hagedorn temperature [4] (see also e.g. [5]). This raises the question whether it is possible to observe, via the conjectured near-horizon-NS5/LST duality, Hagedorn behavior of LST from the thermodynamics of near-horizon NS5-branes.

Indeed, in the work of [6, 7] it was shown that the leading order thermodynamics of non-extremal near-horizon NS5-branes corresponds to the leading order Hagedorn behavior, since one finds the relation $E = TS$ where T is constant. However, since the energy dependence of the entropy near the Hagedorn temperature for a (six-dimensional supersymmetric) string theory is of the form

$$S(E) = \beta_{\mathrm{hg}} E + k \log E \ , \qquad \beta_{\mathrm{hg}} = T_{\mathrm{hg}}^{-1} \tag{1}$$

this apparently only reproduces the first term in the entropy expression. This raises the question of how the NS5-brane background can reproduce (1), or equivalently

$$S(T) = k \frac{T_{\mathrm{hg}}}{T_{\mathrm{hg}} - T} \ , \tag{2}$$

which exhibits the near Hagedorn behavior of the entropy as a function of the temperature. Another related issue is that the leading order NS5-brane thermodynamics has a degenerate phase space. In this talk, we will show that the resolution to these problems is to incorporate string corrections to the NS5-brane supergravity solution [8] (see also [9]). In particular, we reproduce the temperature dependence of (2), where k is a constant determined by the one-loop corrected near-horizon NS5-brane background. This is therefore an example of agreement between statistical thermodynamics of a non-gravitational theory on the one hand and Bekenstein-Hawking thermodynamics of a black brane on the other hand. We note that we can compute the thermodynamics from both sides since we are in a regime in which both LST and the bulk string theory description are weakly coupled. This successful comparison lends further confidence to the conjecture that the various dualities between string and M-theory on near-horizon brane backgrounds and certain non-gravitational theories hold beyond the leading order supergravity solution.

The plan and summary of the talk is as follows. We first review the definition of LST and discuss the dual supergravity solution and its thermodynam-

ics, followed by an analysis of the tree-level and one-loop string corrections arising from the R^4 term in the type II action. We will then show that the latter corrections indeed reproduce the expected temperature dependence (2) of the entropy. Finally, we will comment on the thermodynamics of a number of related non-gravitational theories in six dimensions, known as ODp-brane theories [10, 11], that have recently have been found, along with their supergravity duals. For these theories we derive the thermodynamics as well, showing a different Hagedorn behavior, with critical exponent $-\frac{2}{3}$, the explanation of which is still an open problem.

2   Hagedorn behavior of LST from string corrections to NS5-branes

Six-dimensional LSTs with 16 supercharges can be defined from N coincident NS5-branes in the limit [2]

$$g_s \to 0 \ , \qquad l_s = \text{fixed} \ , \tag{3}$$

with g_s and l_s being respectively the string coupling and string length of type II string theory. In the limit (3) the bulk modes decouple and give rise to non-gravitational theories: The $(1,1)$ LST for type IIB NS5-branes and the $(2,0)$ LST for type IIA NS5-branes, both of type A_{N-1}. Here, , we will use the conjecture that string theory in the background of N NS5-branes is dual to LST [1], in parallel with the AdS/CFT correspondence [12]. In order to be at finite temperature the correspondence for the near-extremal rather than the extremal case will be considered here. We start by considering the supergravity solution of N coincident NS5-branes in type II string theory with r the transverse radius and r_0 the horizon radius

$$ds^2 = H^{-1/4} \left[\left(1 - \frac{r_0^2}{r^2}\right)dt^2 + \sum_{i=1}^{5}(dy^i)^2 + H\left(\left(1 - \frac{r_0^2}{r^2}\right)^{-1}dr^2 + r^2 d\Omega_3^2\right) \right] \ , \tag{4}$$

$$e^\phi = g_s H^{1/2} \qquad , \qquad H = 1 + \frac{N l_s^2}{r^2} \ . \tag{5}$$

Then, if we take the decoupling limit (3) keeping fixed the energy scales

$$u = \frac{r}{g_s l_s^2} \ , \qquad u_0 = \frac{r_0}{g_s l_s^2} \tag{6}$$

one obtains the Einstein-frame metric and dilaton

$$\frac{ds^2}{\sqrt{g_s}l_s} = \sqrt{\frac{u}{\sqrt{N}l_s}}\left[-\left(1-\frac{u_0^2}{u^2}\right)dt^2 + \sum_{i=1}^{5}(dy^i)^2 + Nl_s^2\left(\frac{du^2}{u^2-u_0^2}+d\Omega_3^2\right)\right] ,$$
(7)

$$g_s e^{\phi} = \frac{\sqrt{N}}{l_s u} .$$
(8)

This supergravity solution is conjectured to be dual to a LST [1]. The variable u defined by (6) is kept finite since in type IIB it corresponds to the mass of an open D-string stretching between two NS5-branes with distance r. In type IIA, u/l_s is instead the induced string tension of an open D2-brane stretching between two NS5-branes.

The curvature $e^{-\phi/2}R$ in units of l_s^{-2} and the effective string coupling squared at the horizon radius are respectively of order

$$\varepsilon_D = \frac{1}{N} , \qquad \varepsilon_L = \frac{N}{l_s^2 u_0^2} .$$
(9)

In order for the near-horizon NS5-brane solution (7), (8) to be a dual description of LST we need that $\varepsilon_D \ll 1$ and $\varepsilon_L \ll 1$ which is fulfilled for [7, 1]

$$N \gg 1 , \qquad l_s^2 u_0^2 \gg N ,$$
(10)

so that near-horizon NS5-branes describe LST in the UV-region. The thermo-dynamics[2] of the near-horizon NS5-brane solution is [7]

$$T = \frac{1}{2\pi\sqrt{N}l_s} , \qquad S = \frac{\sqrt{N}V_5}{(2\pi)^4 l_s^3}u_0^2 , \qquad E = \frac{V_5}{(2\pi)^5 l_s^4}u_0^2 , \qquad F = 0 .$$
(11)

Indeed, this is [6, 7], the zeroth order approximation the thermodynamics $S = \beta_{\mathrm{hg}}E$ of a string theory at high temperature by identifying the Hagedorn temperature as

$$T_{\mathrm{hg}} = (2\pi\sqrt{N}l_s)^{-1}$$
(12)

On the other hand, the Hagedorn temperature of a closed string theory with string length $\hat{l}_s$ and central charge c is given by $T_{\mathrm{hg}} = (2\pi\hat{l}_s\sqrt{6/c})^{-1}$ so that for a 5+1 dimensional supersymmetric string theory with $c = 6$, one has $T_{\mathrm{hg}} = (2\pi\hat{l}_s)^{-1}$. Comparing this with (12), on observes that $\hat{l}_s = \sqrt{N}l_s$ is the string length associated with the Hagedorn exponential growth of string states

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[2]The thermodynamics of spinning near-horizon NS5-branes is found in [13, 14].

in LST of type A_{N-1}. In terms of the string tension, $\hat{\tau} = \tau/N$ this shows that the string tension associated with the Hagedorn behavior is quantized in a unit that is a fraction of the ordinary string tension. Indeed, by considering LST as the decoupling limit (3) of type II string theory on an A_{N-1} singularity [15], it is known from the study of string theory on this orbifold singularity [16] that there exist fractional $(1/N)$ branes. Another way to argue that the Hagedorn temperature picks the $1/\hat{l}_s$ string scale is that this is the first one reached when raising the temperature [1].

We now want to include string corrections in the thermodynamics[3], arising from the higher derivative terms and string loops. It follows from (9) that the derivative expansion, which is an expansion in $\alpha' = l_s^2$, becomes an expansion in ε_D for the near-horizon NS5-brane solution. The tree-level expansion of the temperature is

$$T = \frac{1}{2\pi\sqrt{N}l_s}\left(1 + \sum_{i=3}^{\infty} a_i \frac{1}{N^i}\right) , \tag{13}$$

where we used that the leading order correction comes from the R^4 term. This seems to indicate that the Hagedorn temperature is given by (12) only when N is large whereas for finite N it has a different N-dependence. Nevertheless, we expect that there are no tree-level corrections to the Hagedorn temperature so that $a_i = 0$ for all i. This is easily seen to be necessary from the requirement that (13) be equal to (12), since from (13) it would follow that the string tension $\hat{\tau}$ associated with Hagedorn behavior of LST would have corrections $\hat{\tau} = N^{-1}\tau(1 + \mathcal{O}(N^{-1}))$ which clearly cannot make sense as a fractional string tension This fact also follows by noticing that the near-horizon black NS5-brane solution is described by the $SL(2,R)/SO(1,1)\times SU(2)$ exact CFT [7, 18] (see [19] for an exact model of the extremal NS5-brane), which implies that the metric in the string frame, and hence the temperature, does not have tree-level corrections [20].

Turning, to the string loop expansion, this becomes an expansion in ε_L given in (9). In particular, the leading correction to the temperature is therefore generated by the one-loop R^4 term of order $\varepsilon_D^3\varepsilon_L$. Thus, using (9), we can write

$$T = T_{\text{hg}}\left(1 - \frac{\pi^2 b}{2}\frac{1}{N^2 l_s^2 u_0^2}\right) , \tag{14}$$

where T_{hg} is given by (12) and the constant b is a rational number which depends on the exact form of the perturbed solution. Since we assume that

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[3]We refer to [17] for computational details and a general analysis of string corrections to near-horizon brane thermodynamics.

all tree-level corrections to the temperature vanish, (14) is valid for all $u_0 \gg \sqrt{N} l_s^{-1}$. Note that the one-loop correction, has generated a u_0 dependence in the temperature, so that the degenerate phase-space, with constant temperature, now expands and determines the leading order behavior of the thermodynamics. On general grounds, it is expected that the temperature T in (14) is an increasing function of the energy, which means that $b > 0$ in (14). Assuming $b > 0$ we see that (14) strongly suggest that the temperature T_{hg} is a limiting Hagedorn temperature of LST, in the sense that the temperature T of the system can come arbitrarily close to T_{hg} for higher and higher energy $E \propto u_0^2$, but that it never can pass T_{hg}.

Solving (14) for u_0^2 we obtain from (11) the leading order expressions for the entropy as a function of temperature,

$$S(T) = \pi^3 b N \hat{V}_5 \frac{T_{\text{hg}}}{T_{\text{hg}} - T} \ , \tag{15}$$

with $\hat{V}_5 = V_5/(2\pi \hat{l}_s)^5$. We are now in position to compare with the statistical thermodynamics of a 5+1 dimensional closed supersymmetric string theory. To this end we note that from (11) and (10) it follows that $\hat{l}_s E \gg (\hat{V}_5)^{1/5}$ so that the five spatial world volume dimensions are effectively compact, and hence the expected temperature form of the entropy is given by (2). It then follows that the entropy (15) of the one-loop corrected NS5-brane solution and the entropy (2) derived by statistical mechanics for closed strings agree in that they exhibit identical temperature dependence. Moreover, this means that the constant of proportionality in (2) for LST with 16 supercharges is predicted to be

$$k = \pi^3 b N \hat{V}_5 \ . \tag{16}$$

However, there is a remaining puzzle, which has been addressed in [9]. While for the free string, the energy in the compact case is intensive due to the dominance of long strings, for LST the supergravity dual predicts an extensive energy. This hints at a new universality class of interacting strings with a strong self-attractive potential, so that long strings are suppressed and the strings prefer to be coiled [9].

3 Hagedorn behavior of ODp-theories

Recently, new non-gravitational six-dimensional theories have been constructed by turning on a critical electric field on the NS5-brane [10, 11]. These theories

are called Open Dp-brane (ODp) theories[4], and for a given $p = 0, ..., 5$, they contain light open Dp-brane excitations. They are obtained from the Dp-NS5 brane bound state in the limit

$$l_s \to 0 \quad , \quad g_s = \tilde{g}\varepsilon^{\frac{3-p}{4}} \quad , \quad \varepsilon = l_s^4/l_{\text{eff}}^4 \qquad (17)$$

keeping $\tilde{g}$ and l_{eff} finite. The light open p-brane has tension $\frac{1}{2}T_{ODp}$ where $T_{ODp} = ((2\pi)^p \tilde{g} l_{\text{eff}}^{p+1})^{-1}$. The theory obtained in this limit also contains closed strings of tension $1/(2\pi l_{\text{eff}}^2)$ [10, 11]. To see that this is the correct tension, one may compute the tension of a closed string soliton in the low-energy SYM theory, which gives $(2\pi)^2/g_{\text{YM}}^2 = 1/(2\pi l_{\text{eff}}^2)$ since $g_{\text{YM}}^2 = (2\pi)^3 l_{\text{eff}}^2$. Therefore, l_{eff} is the string length associated with "little" closed strings.

The supergravity solution for the Dp-NS5 brane bound state is

$$\begin{aligned} ds^2 &= D^{-1/2}\Big[D\big(-fdt^2 + (dx^1)^2 + \cdots + (dx^p)^2\big) \\ &\quad + (dx^{p+1})^2 + \cdots + (dx^5)^2 + H\big(f^{-1}dr^2 + r^2 d\Omega_3^2\big)\Big] , \end{aligned} \qquad (18)$$

$$e^{2\phi} = HD^{\frac{p-3}{2}} , \qquad (19)$$

$$A_{01\cdots p} = (-1)^p \sin\hat{\theta} \coth\hat{\alpha} D H^{-1} \quad , \quad A_{(p+1)\cdots 5} = (-1)^p \tan\hat{\theta} H^{-1} , \qquad (20)$$

$$f = 1 - \frac{r_0^2}{r^2} \quad , \quad H = 1 + \frac{r_0^2 \sinh^2\alpha}{r^2} \quad , \quad D^{-1} = \cosh^2\theta - \sinh^2\theta H^{-1} , \qquad (21)$$

$$\sinh^2\alpha = \cos^2\hat{\theta} \sinh^2\hat{\alpha} \quad , \quad \cosh^2\theta = \frac{1}{\cos^2\hat{\theta}} . \qquad (22)$$

The near-horizon limit is given by the limit (17) along with

$$\cosh\theta = \frac{1}{\sqrt{\varepsilon}} \quad , \quad r = \tilde{r}\sqrt{\varepsilon} \quad , \quad r_0 = \tilde{r}_0\sqrt{\varepsilon} \quad , \quad x^i = \tilde{x}^i\sqrt{\varepsilon} , \quad i = p+1, ..., 5 \quad (23)$$

from which we obtain the near-horizon solution[5]

$$\begin{aligned} ds^2 &= H^{-1/2}\frac{R}{\tilde{r}}\Bigg[H\frac{\tilde{r}^2}{R^2}\big(-fdt^2 + (dx^1)^2 + \cdots + (dx^p)^2\big) \\ &\quad + (d\tilde{x}^{p+1})^2 + \cdots + (d\tilde{x}^5)^2 + H\big(f^{-1}d\tilde{r}^2 + \tilde{r}^2 d\Omega_3^2\big)\Bigg] , \end{aligned} \qquad (24)$$

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[4]The OD1 and OD2-theories are called $(1, 1)$ and $(2, 0)$ Open Brane Little String Theories (OBLSTs) in [11].

[5]The non-extremal Dp-NS5 solution and its near-horizon limit were given previously in [11] for $p = 1, 2$. The extremal Dp-NS5 solution and its near-horizon limit were given in [21]. For $p = 2, 3$ it was also given in [22].

$$g_s^2 e^{2\phi} = \tilde{g}^2 \frac{l_{\text{eff}}^2 N}{\tilde{r}^2} \left(1 + \frac{\tilde{r}^2}{l_{\text{eff}}^2 N} \right)^{\frac{p-1}{2}} , \tag{25}$$

$$T_{Dp} A_{01\cdots p} = (-1)^p T_{ODp} \frac{\tilde{r}^2}{l_{\text{eff}}^2 N} \quad , \quad T_{D(4-p)} A_{(p+1)\cdots 5} = (-1)^p T_{OD(4-p)} H^{-1} , \tag{26}$$

$$f = 1 - \frac{\tilde{r}_0^2}{\tilde{r}^2} \quad , \quad H = 1 + \frac{l_{\text{eff}}^2 N}{\tilde{r}^2} , \tag{27}$$

with $T_{Dp} = ((2\pi)^p g_s l_s^{p+1})^{-1}$. The curvature in units of l_s^{-2} is of order $\mathcal{C} = \left(N^2 + \tilde{r}^2 N l_{\text{eff}}^{-2} \right)^{-1/2}$ and we need $\mathcal{C} \ll 1$ and $g_s e^\phi \ll 1$ in order for the near horizon solution (24)-(27) to describe ODp-theory. We see that the solution (24)-(27) has two phases, the NS5-brane phase for $\tilde{r} \ll l_{\text{eff}} \sqrt{N}$ and the delocalized Dp-brane phase for $\tilde{r} \gg l_{\text{eff}} \sqrt{N}$. In the NS5-brane phase the theory is the ordinary $(1,1)$ LST or $(2,0)$ LST which reduces to SYM$_{5+1}$ or $(2,0)$ SCFT at low energies. The Dp-brane phase is instead a theory of closed "little" strings on a special geometry.

The leading order thermodynamics is[6]

$$T = \frac{1}{2\pi l_{\text{eff}} \sqrt{N}} \quad , \quad S = l_{\text{eff}} \sqrt{N} \frac{\tilde{V}_5}{\tilde{g}^2 (2\pi)^4} \tilde{r}_0^2 \quad , \quad F = 0 \quad , \quad E = \frac{\tilde{V}_5}{(2\pi)^5} \tilde{r}_0^2 . \tag{28}$$

The curvature at the horizon radius in units of l_s^{-2} and the effective string coupling squared at the horizon are respectively of order

$$\varepsilon_D = \left(N^2 + \tilde{r}_0^2 N l_{\text{eff}}^{-2} \right)^{-1/2} \quad , \qquad \varepsilon_L = \tilde{g}^2 \frac{l_{\text{eff}}^2 N}{\tilde{r}_0^2} \left(1 + \frac{\tilde{r}_0^2}{l_{\text{eff}}^2 N} \right)^{\frac{p-1}{2}} . \tag{29}$$

The leading order thermodynamics (28) is valid for $\varepsilon_D \ll 1$ and $\varepsilon_L \ll 1$. As for ordinary LST, the thermodynamics (28) exhibits leading order Hagedorn behavior with Hagedorn temperature

$$T_{\text{hg}} = (2\pi l_{\text{eff}} \sqrt{N})^{-1} . \tag{30}$$

We see that this temperature is the Hagedorn temperature associated with the "little" closed strings in ODp-theory [11]. In fact, when the supergravity dual (24)-(27) is valid we have that $T \sim T_{\text{hg}}$, which we take to mean that the "little" closed strings are the fundamental degrees of freedom in this region of phase space. We note that the Hagedorn temperature (30) was also found using different arguments in [23].

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[6]This has previously been computed for the OD1 and OD2-theories in [11].

We now want to compute the leading correction to the temperature, as done above for ordinary LST, since this gives the temperature dependence of the entropy. Since we want to consider the high energy behavior of the thermodynamics we assume that we are in the Dp-brane phase $\tilde{r}_0 \gg l_{\text{eff}}\sqrt{N}$. Contrary to the NS5-brane case, the leading correction in this case comes from the leading tree-level R^4 term at order ε_D^3. We therefore find

$$T = T_{\text{hg}}\left(1 - a^3 \frac{l_{\text{eff}}^3}{N^{3/2}\tilde{r}_0^3}\right) \,, \tag{31}$$

where a is an undetermined constant. The resulting leading order expression for the entropy as a function of temperature is

$$S(T) = a^2 \frac{l_{\text{eff}}^3}{\sqrt{N}} \frac{\tilde{V}_5}{\tilde{g}^2(2\pi)^4}\left(\frac{T_{\text{hg}}}{T_{\text{hg}} - T}\right)^{2/3} \,. \tag{32}$$

Comparing (32) with the entropy (15) for ordinary LST we see that the critical behavior for high energies is different. For ordinary LST the entropy is proportional to $(T_{\text{hg}} - T)^{-1}$, while for the ODp-theories the entropy is proportional to $(T_{\text{hg}} - T)^{-2/3}$, as already noticed in [8, 24, 11]. Even though the high energy behavior of ODp-theory is dominated by "little" closed strings just like for ordinary LST, we interpret this difference as arising from the fact that the "little" closed strings of ODp-theory live on a different kind of geometry, e.g. for $p = 1$ the space-time is space-time non-commutative. We call this phase the deformed LST phase. It would be very interesting to reproduce the critical exponent $-\frac{2}{3}$ of the entropy from the statistical mechanics of closed strings on one of the ODp-theory geometries.

References

[1] O. Aharony, M. Berkooz, D. Kutasov, and N. Seiberg, "Linear dilatons, NS5-branes and holography," *JHEP* **10** (1998) 004, `hep-th/9808149`.

[2] N. Seiberg, "New theories in six-dimensions and matrix description of m theory on T^5 and T^5/Z_2," *Phys. Lett.* **B408** (1997) 98–104, `hep-th/9705221`; M. Berkooz, M. Rozali, and N. Seiberg, "Matrix

description of M theory on T^4 and T^5," *Phys. Lett.* **B408** (1997) 105–110, `hep-th/9704089`.

[3] R. Dijkgraaf, E. Verlinde, and H. Verlinde, "Notes on matrix and micro strings," *Nucl. Phys. Proc. Suppl.* **62** (1998) 348–362, `hep-th/9709107`; A. Losev, G. Moore, and S. L. Shatashvili, "M & m's," *Nucl. Phys.* **B522** (1998) 105–124, `hep-th/9707250`.

[4] R. Hagedorn, "Statistical thermodynamics of strong interactions at high- energies," *Nuovo Cim. Suppl.* **3** (1965) 147–186.

[5] P. Salomonson and B.-S. Skagerstam, "On superdense superstring gases: A heretic string model approach," *Nucl. Phys.* **B268** (1986) 349; J. J. Atick and E. Witten, "The Hagedorn transition and the number of degrees of freedom of string theory," *Nucl. Phys.* **B310** (1988) 291; R. Brandenberger and C. Vafa, "Superstrings in the early universe," *Nucl. Phys.* **B316** (1989) 391.

[6] J. M. Maldacena, "Statistical entropy of near extremal five-branes," *Nucl. Phys.* **B477** (1996) 168–174, `hep-th/9605016`.

[7] J. M. Maldacena and A. Strominger, "Semiclassical decay of near extremal five-branes," `hep-th/9710014`.

[8] T. Harmark and N. A. Obers, "Hagedorn behaviour of little string theory from string corrections to NS5-branes," *Phys. Lett.* **B485** (2000) 285, `hep-th/0005021`.

[9] M. Berkooz and M. Rozali, "Near Hagedorn dynamics of NS fivebranes, or a new universality class of coiled strings," *JHEP* **05** (2000) 040, `hep-th/0005047`.

[10] R. Gopakumar, S. Minwalla, N. Seiberg, and A. Strominger, "OM theory in diverse dimensions," *JHEP* **08** (2000) 008, `hep-th/0006062`.

[11] T. Harmark, "Open branes in space-time non-commutative little string theory," `hep-th/0007147`.

[12] J. Maldacena, "The large N limit of superconformal field theories and supergravity," *Adv. Theor. Math. Phys.* **2** (1998) 231–252, `hep-th/9711200`; O. Aharony, S. S. Gubser, J. Maldacena, H. Ooguri, and Y. Oz, "Large N field theories, string theory and gravity," *Phys. Rept.* **323** (2000) 183, `hep-th/9905111`.

[13] K. Sfetsos, "Rotating NS5-brane solution and its exact string theoretical description," *Fortsch. Phys.* **48** (2000) 199–204, `hep-th/9903201`.

[14] T. Harmark and N. A. Obers, "Thermodynamics of spinning branes and their dual field theories," *JHEP* **01** (2000) 008, `hep-th/9910036`; T. Harmark and N. A. Obers, "Thermodynamics of field theories from spinning branes," `hep-th/0002250`.

[15] H. Ooguri and C. Vafa, "Two-dimensional black hole and singularities of CY manifolds," *Nucl. Phys.* **B463** (1996) 55–72, `hep-th/9511164`.

[16] M. R. Douglas and G. Moore, "D-branes, quivers, and ALE instantons," `hep-th/9603167`; C. V. Johnson and R. C. Myers, "Aspects of type IIB theory on ALE spaces," *Phys. Rev.* **D55** (1997) 6382–6393, `hep-th/9610140`.

[17] J. Correia, T. Harmark, and N. A. Obers. In preparation.

[18] K. Sfetsos, "On (multi-)center branes and exact string vacua," `hep-th/9812165`. In *Corfu 1998, Quantum aspects of gauge theories, supersymmetry and unification* 303-312.

[19] C. G. Callan Jr., J. A. Harvey, and A. Strominger, "World sheet approach to heterotic instantons and solitons," *Nucl. Phys.* **B359** (1991) 611; C. G. Callan Jr., J. A. Harvey, and A. Strominger, "Worldbrane actions for string solitons," *Nucl. Phys.* **B367** (1991) 60–82.

[20] A. A. Tseytlin, "On field redefinitions and exact solutions in string theory," *Phys. Lett.* **B317** (1993) 559–564, `hep-th/9308042`.

[21] M. Alishahiha, Y. Oz, and J. G. Russo, "Supergravity and light-like non-commutativity," `hep-th/0007215`.

[22] M. Alishahiha, Y. Oz, and M. M. Sheikh-Jabbari, "Supergravity and large N noncommutative field theories," *JHEP* **11** (1999) 007, `hep-th/9909215`; M. Alishahiha, "On type II NS5-branes in the presence of an RR field," `hep-th/0002198`.

[23] S. S. Gubser, S. Gukov, I. R. Klebanov, M. Rangamani, and E. Witten, "The Hagedorn transition in non-commutative open string theory," `hep-th/0009140`.

[24] T. Harmark, "Supergravity and space-time non-commutative open string theory," *JHEP* **07** (2000) 043, `hep-th/0006023`.